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Magnetic moments of the exotic pentaquark baryons within the chiral quark-soliton model*

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We present in this talk recent results of the magnetic moments of the baryon antidecuplet within the framework of the chiral quark-soliton model in the chiral limit. The dynamic parameters of the model are fixed by using the experimental data for those of the baryon octet. Sum rules for the magnetic moments are derived. We found that the magnetic moments of the baryon antidecuplet have opposite signs to their charges. The magnetic moments of the neutral baryon antidecuplet turn out to be compatible with zero.

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I. INTRODUCTION

The LEPS collaboration [1] announced the finding of the Θ^+ consisting of four quarks and one anti-quark (uudd$\bar{\rm s}$), motivated by a theoretical prediction of the chiral soliton model [2]. Since then, a great amount of experimental and theoretical works [3, 4, 5, 6, 7, 8, 9, 10, 11, 12, 13, 14, 15, 16, 17, 18, 19, 20, 21, 22, 23] has been published.

In order to describe the photo-production of pentaquark baryons [24, 25, 26], we need information on their magnetic moments. Since there is no experimental data available, one needs a theoretical guideline to estimate them. Recently, the present authors investigated the magnetic moments of the pentaquark baryons within the chiral quark-soliton model in a "*model-independent*" way [27, 28, 29]. In this approach the dynamical model parameters are fixed by using the experimental data of the octet magnetic moments [30]. However, not all parameters can be constrained that way. Hence Ref. [30] used some additional information based on the dynamical model calculations. Therefore the analysis of Ref. [30] was not self-consistent. In this talk, we will present the recent results for the magnetic moments of the baryon antidecuplet in the chiral limit with the complete set of parameters fixed from the experimental data.

II. FORMALISM

The magnetic moments of the baryon antidecuplet can be defined as the following one-current baryon matrix element:

$$\langle B_{\overline{10}}|\bar{\psi}(z)\gamma_\mu\hat{Q}\psi(z)|B_{\overline{10}}\rangle, \tag{1}$$

where $\hat{Q}$ denotes the charge operator of quarks in SU(3) flavor space, defined by

$$\hat{Q} = \begin{pmatrix} \frac{2}{3} & 0 & 0 \\ 0 & -\frac{1}{3} & 0 \\ 0 & 0 & -\frac{1}{3} \end{pmatrix} = \frac{1}{2}\left(\lambda^3 + \frac{1}{\sqrt{3}}\lambda^8\right). \tag{2}$$

In the nonrelativistic limit, the Sachs form factors G_E and G_M can be related to the time and space components of the $\mathrm{U}_V(3)$ vector currents, respectively:

$$\langle B_{\overline{10}}(p')|\bar{\psi}(z)\gamma_0\hat{Q}\psi(z)|B_{\overline{10}}(p)\rangle = G_E^{B_{\overline{10}}}(Q^2), \tag{3}$$

$$\begin{aligned}&\langle B_{\overline{10}}(p')|\bar{\psi}(z)\gamma_i\hat{Q}\psi(z)|B_{\overline{10}}(p)\rangle \\ &= \frac{1}{2M_N}G_M^{B_{\overline{10}}}(Q^2)i\epsilon_{ijk}q^j\langle s'|\sigma_k|s\rangle,\end{aligned} \tag{4}$$

where σ_k denotes Pauli spin matrices while $|s\rangle$ is the corresponding spin state of the baryon. The magnetic moments $\mu_{B_{\overline{10}}}$ corresponding to the vector currents are identified with $G_M^{B_{\overline{10}}}(0)$.

The collective magnetic moment operator can be obtained schematically by differentiating the effective chiral action with the external source corresponding to the magnetic moments as follows:

$$\hat{\mu}_k = \frac{\delta}{\delta s_k}S_{\rm eff}$$

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$$= -N_c \frac{\delta}{\delta s_k} \mathrm{Tr}\log\left[i\partial_4 + iH(U_c^{\gamma_5}) - \Omega + i\gamma_4 R^\dagger \hat{m} R \right.$$
$$\left. -is_m \epsilon_{ilm}\gamma_4\gamma_i x_l R^\dagger \hat{Q} R\right], \quad (5)$$

where $H(U^{\gamma_5})$ is the one-body Dirac Hamiltonian defined by

$$H(U^{\gamma_5}) = \frac{\boldsymbol{\alpha}\cdot\boldsymbol{\nabla}}{i} + \beta M U^{\gamma_5}. \quad (6)$$

U^{γ_5} stands for the chiral soliton field:

$$U^{\gamma_5} = \frac{1+\gamma_5}{2}U + \frac{1-\gamma_5}{2}U^\dagger \quad (7)$$

with the trivial embedding

$$U_c = \begin{pmatrix} U_{\mathrm{SU(2)}} & 0 \\ 0 & 1 \end{pmatrix}. \quad (8)$$

Here, $U(SU(2))$ denotes the SU(2) soliton field. $\Omega = \frac{1}{2}\lambda_a\Omega_a$ in Eq.(5) designates the angular velocity of the soliton, which is related to the right angular momentum operator after the zero-mode quantization:

$$\mathcal{R}_a = -\Omega_b I_{ba} + \frac{N_c}{2\sqrt{3}}\delta_{8a} - 2K_{ab} m_8 D_{8b}^{(8)}(R) \quad (9)$$

with the current quark masses

$$\hat{m} = \begin{pmatrix} m_{\mathrm{u}} & 0 & 0 \\ 0 & m_{\mathrm{d}} & 0 \\ 0 & 0 & m_{\mathrm{s}} \end{pmatrix} = m_0\mathbf{1} + m_8\lambda_8. \quad (10)$$

I_{ab} and K_{ab} stand for the moments of inertia. Taking into account the rotational $1/N_c$ corrections as well as the linear m_{s} corrections, we arrive at the following collective operator of the magnetic moments:

$$\begin{aligned}\hat{\mu}_3 &= w_1\, D_{Q3}^{(8)} + w_2 d_{pq3} D_{Qp}^{(8)}\hat{S}_q + \frac{w_3}{\sqrt{3}} D_{Q8}^{(8)}\hat{S}_3 \\ &+ m_{\mathrm{s}}\left[\frac{w_4}{\sqrt{3}} d_{pq3} D_{Qp}^{(8)} D_{8q}^{(8)} \right. \\ &+ w_5\left(D_{Q3}^{(8)}D_{88}^{(8)} + D_{Q8}^{(8)}D_{83}^{(8)}\right) \\ &+ \left. w_6\left(D_{Q3}^{(8)}D_{88}^{(8)} - D_{Q8}^{(8)}D_{83}^{(8)}\right)\right]. \quad (11)\end{aligned}$$

where the dynamical variables w_i contain information of the dynamics of the chiral soliton, which are independent of baryons considered. They can be generically written in terms of the inertia parameters of the soliton in the χQSM:

$$\sum_{m,n}\langle n|\Gamma_1|m\rangle\langle m|\Gamma_2|n\rangle \mathcal{R}(E_n, E_m, \Lambda), \quad (12)$$

where Γ_i denote spin-isospin operators acting on the quark eigenstates $|n\rangle$ of the one-body Dirac Hamiltonian (6) in the soliton-background field. The double sum over all the eigenstates can be evaluated numerically [31, 32, 33]. Since its sea part diverges, we need the regularization expressed by $\mathcal{R}$ with the cut-off parameter Λ fixed to the pion decay constant. In this work, we will not calculate the dynamical variables w_i numerically but we will constrain them using the experimental data of the octet magnetic moments [28, 29]. $D_{ab}^{(\mathcal{R})}(R)$ stands for the SU(3) Wigner function, $R(t)$ is the time-dependent SU(3) matrix responsible for the rotation of the soliton in the collective coordinate space [31, 34]. $\hat{J}_a$ denotes an operator of the generalized spin acting on the baryonic wave functions $\psi_{B_\mathcal{R}}(R)$. Since the SU(3) symmetry breaking introduce the mixing of the pure antidecuplet states with higher representations, we have to calculate the wave-function corrections. However, we will present here only the results of the magnetic moments in the chiral limit, namely, we will put $m_{\mathrm{s}} = 0$ in Eq.(5).

In order to evaluate the magnetic moments of the baryon antidecuplet, we have to sandwich the collective opreator in Eq.(5) between the collective baryon $\overline{\mathbf{10}}$ states:

$$\mu_{B_{\overline{10}}} = \int dR \psi^*_{B_{\overline{10}}}(R)\hat{\mu}(R)\psi_{B_{\overline{10}}}(R), \quad (13)$$

where the collective wave functions $\psi_{B_\mathcal{R}}(R)$ are defined as follows:

$$\psi_{B_\mathcal{R}}(R) = \sqrt{\dim(\mathcal{R})}(-1)^{J_3 - Y'/2} D^{(\mathcal{R})*}_{Y,T,T_3;Y',J,-J_3}(R). \quad (14)$$

Here $\mathcal{R}$ stands for the allowed irreducible representations of the SU(3) flavor group, *i.e.* $\mathcal{R} = 8, 10, \overline{10}, \cdots$ and Y, T, T_3 are the corresponding hypercharge, isospin, and its third component, respectively. Right hypercharge Y' is constrained to be unity for the physical spin states for which J and J_3 are spin and its third component. Note that under the action of left (flavor) generators $\hat{T}_\alpha = -D^{(8)}_{\alpha\beta}\hat{J}_\beta$ $\psi_{B_\mathcal{R}}$ transforms like a tensor in representation $\mathcal{R}$, while under the right generators $\hat{J}_\alpha$ like a tensor in $\mathcal{R}^*$ rather than $\mathcal{R}$. This is the reason why operators like the one multiplied by w_2 in Eq.(5) have different matrix elements for the decuplet (which is spin 3/2) and antidecuplet (which is spin 1/2). The other two operators multiplied by $w_{1,3}$ have the same matrix elements between decuplet and antidecuplet states.

The matrix elements of Eq.(13) are expressed in terms of SU(3) Clebsch-Gordan coefficients [35]. Having scrutinized the results, we find the following simple expression:

$$\mu_{B_{\overline{10}}} = -\frac{1}{12}\left(w_1 + \frac{5}{2}w_2 - \frac{1}{2}w_3\right) Q_{B_{\overline{10}}}\, J_3, , \quad (15)$$
$$\mu_{B_{10}} = -\frac{1}{12}\left(w_1 - \frac{1}{2}w_2 - \frac{1}{2}w_3\right) Q_{B_{10}}\, J_3, \quad (16)$$

where $Q_{B_{\overline{10}}}$ is the charge of the antidecuplet expressed by the Gell-Mann–Nishijima relation:

$$Q_{B_{\overline{10}}} = T_3 + \frac{Y}{2}. \quad (17)$$

J_3 is the corresponding third component of the spin.

III. NUMERICAL FITS

In order to fit the parameters w_i, it is convenient to introduce two parameters consisting of w_1, w_2 and w_3:

$$v = \frac{1}{60}\left(w_1 - \frac{1}{2}w_2\right), \quad w = \frac{1}{120}\, w_3. \tag{18}$$

In Ref. [28] the octet and decuplet magnetic moments were expressed as follows:

$$\begin{aligned}
\mu_p &= \mu_{\Sigma^+} = -8v + 4w, \\
\mu_n &= \mu_{\Xi^0} = 6v + 2w, \\
\mu_\Lambda &= -\mu_{\Sigma^0} = 3v + w, \\
\mu_{\Sigma^-} &= \mu_{\Xi^-} = 2v - 6w, \\
\mu_{B_{10}} &= \frac{15}{2}\left(-v + w\right) Q_{B_{10}}.
\end{aligned} \tag{19}$$

which are in fact the well-known SU(3) formulae for the magnetic moments.

On the other hand, magnetic moments of the baryon antidecuplet (15) can be rewritten as:

$$\mu_{B_{\overline{10}}} = \left[\frac{5}{2}\left(-v + w\right) - \frac{1}{8}w_2\right] Q_{B_{\overline{10}}}. \tag{20}$$

Equation (20) is different from the decuplet in Eq.(16) by the second term proportional to w_2. The factor three difference in the first term between Eq.(19) and Eq.(20) is due to the fact that the baryon antidecuplet has spin 1/2, while the decuplet has 3/2.

Using Eq.(20), we are able to derive the sum rules which are similar to the generalized Coleman and Glashow sum rules [36] in the chiral limit:

$$\begin{aligned}
\mu_{\Sigma^0_{\overline{10}}} &= \frac{1}{2}\left(\mu_{\Sigma^+_{\overline{10}}} + \mu_{\Sigma^-_{\overline{10}}}\right), \\
\mu_{\Xi^+_{3/2}} + \mu_{\Xi^{--}_{3/2}} &= \mu_{\Xi^0_{3/2}} + \mu_{\Xi^-_{3/2}}, \\
\sum \mu_{B_{\overline{10}}} &= 0.
\end{aligned} \tag{21}$$

As discussed in Ref. [28], there are different ways to fix the parameters v and w by using the experimental data of the octet magnetic moments. Here, we simply fit the proton and neutron magnetic moments (fit I):

$$\begin{aligned}
v &= (2\mu_{\rm n} - \mu_{\rm p})/20 &= -0.331, \\
w &= (4\mu_{\rm n} + 3\mu_{\rm p})/20 &= 0.037,
\end{aligned} \tag{22}$$

and use the following "average" values (fit II):

$$\begin{aligned}
v &= \left(2\mu_{\rm n} - \mu_{\rm p} + 3\mu_{\Xi^0} + \mu_{\Xi^-} - 2\mu_{\Sigma^-} - 3\mu_{\Sigma^+}\right)/60 \\
&= -0.268,
\end{aligned} \tag{23}$$

$$\begin{aligned}
w &= \left(3\mu_{\rm p} + 4\mu_{\rm n} + \mu_{\Xi^0} - 3\mu_{\Xi^-} - 4\mu_{\Sigma^-} - \mu_{\Sigma^+}\right)/60 \\
&= 0.063.
\end{aligned} \tag{24}$$

to fix parameters v and w. It was shown in Ref. [28] that combinations of Eq.(24) are independent of the linear corrections due to the nonzero strange quark mass $m_{\rm s}$. Thus, fit II is also valid when the SU(3)-symmetry breaking is taken into account, while fit I will be changed by the corrections of order $\mathcal{O}(m_{\rm s})$.

While v and w can be fixed by Eq.(24), we are not able to get w_2 from the analysis in the chiral limit. In order to fix it, we have to carry out the full analysis with the $m_{\rm s}$ corrections considered in Ref. [37]. Otherwise, we have to take it from the model calculation [30]. The value of w_2 obtained with the SU(3) symmetry breaking depends on the pion-nucleon Σ term [23], and for the values of $\Sigma \sim 70$ Mev we get:

$$m_s\, w_2 = 9.81. \tag{25}$$

Compared to the value from the model $m_s\, w_2^{\chi\rm QSM} \sim 5$ used in Ref. [37], it is almost two times larger.

IV. RESULTS AND DISCUSSION

The results of these fits are listed in Table I. We see

	exp.	fit I	fit II	χQSM
p	2.79	input	2.39	2.27
n	−1.91	input	−1.49	−1.55
Λ	−0.61	−0.96	−0.74	−0.78
Σ^+	2.46	2.79	2.38	2.27
Σ^0	(0.65)	0.96	0.74	0.78
Σ^-	−1.16	−0.89	−0.90	−0.71
Ξ^0	−1.25	−1.91	−1.49	−1.55
Ξ^-	−0.65	−0.89	−0.90	−0.71
Δ^{++}	4.52	5.52	4.92	4.47
Ω^-	−2.02	−2.76	−2.46	−2.23
Θ^+	?	−0.31	−0.40	0.12
Ξ^{--}	?	0.62	0.8	−0.24

TABLE I: The magnetic moments of the baryon octet, decuplet, and antidecuplet in the chiral limit. The experimental value for the Δ^{++} magnetic moments is taken from Ref.[38].

that the quality of these fits is rather poor reaching in its worst case about 25% accuracy, which indicates the importance of the SU(3)-symmetry breaking corrections. The results for the baryon antidecuplet are rather remarkable, because their magnetic moments have opposite signs to their charges. Since in the chiral limit the decuplet and antidecuplet magnetic moments are proportional to their charges, those for the neutral baryons turn out to be zero.

If in the χQSM one artificially sets the soliton size $r_0 \to 0$, then the model reduces to the free valence quarks which, however, "remember" the soliton structure. In this limit, many quantities, for example the axial-vector couplings, are given as ratios of the group-theoretical factors [39]. In the case of magnetic moments the pertinent expressions are given as a product of the group-theoretical factor and the model-dependent integral which we shall in what follows denote by K [40].

Constants $w_{1,2,3}$ entering Eq.(5) are expressed in terms of the inertia parameters in the following way:

$$w_1 = M_0 - \frac{M_1^{(-)}}{I_1^{(+)}}, \quad w_2 = -2\frac{M_2^{(+)}}{I_2^{(+)}}, \quad w_3 = -2\frac{M_1^{(+)}}{I_1^{(+)}}. \tag{26}$$

For the soliton size $r_0 \to 0$ we have [40]:

$$\begin{aligned} M_0 \ &\to \ -2K \ , \quad \frac{M_1^{(-)}}{I_1^{(+)}} \to \frac{4}{3}K, \\ \frac{M_2^{(+)}}{I_2^{(+)}} \ &\to \ -\frac{4}{3}K \qquad \frac{M_1^{(+)}}{I_1^{(+)}} \to -\frac{2}{3}K \ , \end{aligned} \tag{27}$$

which give

$$v = -\frac{7}{90}K, \quad w = \frac{1}{90}K, \quad w_3 = \frac{4}{3}K, \tag{28}$$

yielding the magnetic moments of the proton and neutron as follows:

$$\mu_p = \frac{2}{3}K, \quad \mu_n = -\frac{4}{9}K. \tag{29}$$

Hence, the ratio of the proton magnetic moment to the neutron one takes the value from the nonrelativistic quark model:

$$\frac{\mu_p}{\mu_n} = -\frac{2}{3}. \tag{30}$$

We get for the antidecuplet magnetic moments:

$$\mu_{B_{\overline{10}}} = -\frac{1}{3}KQ_{B_{\overline{10}}} \tag{31}$$

which agrees in sign with the phenomenological value of Table I (note that K is positive in view of Eq.(29)). Extracting K from proton or neutron magnetic moments we get $K = 3.4$ and 4.3 respectively. These values lead to rather large, bur negative, value of $\mu_{\Theta^+} = -1.15 \sim -1.4$ respectively.

V. SUMMARY

In the present talk, we determined the magnetic moments of the positive parity baryon antidecuplet in a "*model independent*" analysis, based on the chiral quark-soliton model in the chiral limit. Starting from the collective operators with dynamical parameters fixed by experimental data, we were able to obtain the magnetic moments of the baryon antidecuplet. The expression for the magnetic moments of the antidecuplet is different from those of the baryon decuplet. We found that the magnetic moment of μ_{Θ^+} is about $-0.3 \sim -0.4\,\mu_N$ which differs from the recent results of Refs. [18, 20, 21] and our previous estimate [30] where generally μ_{Θ^+} is small and positive.

In the present talk, we have presented results in the chiral limit. The SU(3)-symmetry breaking effects will definitely make the magnetic moments of the baryon antidecuplet deviate from those of the present paper. There are two different sources of the SU(3)-symmetry breaking effects: one comes from the collective operator, the other arises from the fact that the collective wave functions of the baryon antidecuplet are mixed with the octet, eikosiheptaplet (**27**), and $\overline{\mathbf{35}}$ representations. Moreover, non-analytical symmetry breaking effects are of importance [41]. The effect of the SU(3)-symmetry breaking on the magnetic moments of the antidecuplet baryons has been studied and the results will soon appear [37].

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[1] T. Nakano *et al.* Phys. Rev. Lett. **91** (2003) 012002.
[2] D. Diakonov, V. Petrov and M. V. Polyakov, Z. Phys. A**359** (1997) 305 [hep-ph/9703373].
[3] M. Praszałowicz, talk at *Workshop on Skyrmions and Anomalies*, M. Jeżabek and M. Praszałowicz eds., World Scientific 1987, page 112 and Phys. Lett. B**575** (2003) 234 [hep-ph/0308114].
[4] S. Stepanyan *et al.* [CLAS Collaboration], hep-ex/0307018.
[5] V. V. Barmin *et al.* [DIANA Collaboration], hep-ex/0304040.
[6] J. Barth [SAPHIR Collaboration], hep-ex/0307083.
[7] A. Aktas *et al.* [H1 Collaboration], hep-ex/0403017.
[8] S. Chekanov *et al.* [ZEUS Collaboration], hep-ex/0403051.
[9] A.E. Asratyan, A.G. Dolgolenko, M.A. Kubantsev,

hep-ex/0309042.

[10] C. Alt *et al.* [NA49 Collaboration], hep-ex/0310014.

[11] R.L. Jaffe and F. Wilczek, hep-ph/0307341;
M. Karliner and H. Lipkin, hep-ph/0307243;
S. Capstick, Ph.R. Page and W. Roberts, Phys. Lett. B**570** (2003) 185 [hep-ph/0307019];
C.E. Carlson, Ch.D. Caronne, H.J. Kwee and V. Nazaryan, Phys. Lett. B**573** (2003) 101 [hep-ph/0307396]; hep-ph/0310038;
E. Shuryak and I. Zahed, hep-ph/0310270.

[12] H. Weigel, Eur. Phys. J. **A2** (1998) 391 [hep-ph/9804260];
H. Walliser and V.B. Kopeliovich, Nucl. Phys. B**660** (2003) 156 [hep-ph/0302013];
D. Diakonov and V. Petrov, hep-ph/0309203.

[13] N. Itzhaki, I.R. Klebanov, P. Ouyang and L. Rastelli, hep-ph/0309305.

[14] F. Csikor, Z. Fodor, S.D. Katz and T.D. Kovacs, hep-lat/0309090;
S. Sasaki, hep-lat/0310014.

[15] D.E. Kahana and S.H. Kahana, hep-ph/0310026;
F. J. Llanes-Estrada, E. Oset and V. Mateu, nucl-th/0311020.

[16] F. Stancu and D. O. Riska, hep-ph/0307010.

[17] J. Sugiyama, T. Doi and M. Oka, hep-ph/0309271.

[18] Y. R. Liu, P. Z. Huang, W. Z. Deng, X. L. Chen and S. L. Zhu, hep-ph/0312074.

[19] S.L. Zhu, hep-ph/0307345.

[20] R. Bijker, M. M. Giannini and E. Santopinto, hep-ph/0403029. P.-Z. Huang, W.-Z. Deng, Xi.-L. Chen, S.L. Zhu, hep-ph/0311108; Q. Zhao, Phys. Rev. D **69**, 053009 (2004) [hep-ph/0310350].

[21] D. K. Hong, Y. J. Sohn and I. Zahed, hep-ph/0403205.

[22] Y. Oh, H. Kim and S. H. Lee, hep-ph/0310117;
D. Diakonov and V. Petrov, hep-ph/0310212;
R. Bijker, M.M. Giannini and E. Santopinto, hep-ph/0310281;
T. Hyodo, A. Hosaka and E. Oset, nucl-th/0307105;
C. Hanhart *et al.*, hep-ph/0312236;
S. I. Nam, A. Hosaka and H.-Ch. Kim, hep-ph/0401074.

[23] M. Praszałowicz, Acta Phys. Polon. B **35**, 1625 (2004) [hep-ph/0402038]; J. R. Ellis, M. Karliner and M. Praszałowicz, JHEP **0405**, 002 (2004) [hep-ph/0401127].

[24] W. Liu and C. M. Ko, nucl-th/0308034.

[25] S. I. Nam, A. Hosaka and H.-Ch. Kim, Phys. Lett. B **579** (2004) 43 [hep-ph/0308313];

[26] Y. s. Oh, H. c. Kim and S. H. Lee, hep-ph/0310019.

[27] G. S. Adkins and C. R. Nappi, Nucl. Phys. B **249** (1985) 507.

[28] H.-Ch. Kim, M. Praszałowicz and K. Goeke, Phys. Rev. **D57** (1998) 2859 [hep-ph/9706531].

[29] H.-Ch. Kim, M. Praszałowicz, M. V. Polyakov and K. Goeke, Phys. Rev. **D58** (1998) 114027 [hep-ph/9801295].

[30] H.-Ch. Kim and M. Praszałowicz, Phys. Lett. B **585** (2004) 99 [hep-ph/0308242].

[31] A. Blotz, D. Diakonov, K. Goeke, N. W. Park, V. Petrov and P. V. Pobylitsa, Nucl. Phys. **A555** (1993) 765.

[32] K. Goeke, A. Z. Gorski, F. Gruemmer, T. Meissner, H. Reinhardt and R. Wunsch, Phys. Lett. B**256** (1991) 321.

[33] M. Wakamatsu and H. Yoshiki, Nucl. Phys. **A524** (1991) 561.

[34] C. V. Christov *et al.*, Prog. Part. Nucl. Phys. **37** (1996) 91. [hep-ph/9604441].

[35] J.J. de Swart, Rev. Mod. Phys. **35** (1973) 916;
T.A. Kaeding and H.T. Williams, Computer Phys. Comm. **98**(1996) 398.

[36] S. R. Coleman and S. L. Glashow, Phys. Rev. Lett. **6** (1961) 423.

[37] G.S. Yang, H.-Ch. Kim, M. Praszałowicz, and K. Goeke, in preparation.

[38] A. Bosshard *et al.*, Phys. Rev. **D44** (1991) 930.

[39] M. Praszałowicz, A. Blotz and K. Goeke, Phys. Lett. B**354** (1995) 415 [hep-ph/9505328].

[40] M. Praszałowicz, T. Watabe and K. Goeke, Nucl. Phys. A**647** (1999) 49 [hep-ph/9806431].

[41] E. Jenkins, M. E. Luke, A.V. Manohar and Martin J. Savage, Phys. Lett. B**302** (1993) 482, Erratum ibid. B**388** (1996) 866 [hep-ph/9212226];
U.-G. Meissner and S. Steininger, Nucl. Phys. B**499** (1997) 349 [hep-ph/9701260].